SOFT SUPERSYMMETRY BREAKING INDUCED BY HIGHER-DERIVATIVE SUPERGRAVITATION IN THE ELECTROWEAK STANDARD MODEL

AHMED HINDAWI, BURT A. OVRUT, AND DANIEL WALDRAM

Department of Physics, University of Pennsylvania
Philadelphia, PA 19104-6396, USA

ABSTRACT. We show how spontaneous supersymmetry breaking in the vacuum state of higher-derivative supergravity is transmitted, as explicit soft supersymmetry-breaking terms, to the effective Lagrangian of the standard electroweak model. The general structure of the soft supersymmetry breaking terms is presented and a new scenario for understanding the gauge hierarchy problem, based on the functional form of these terms, is discussed.

PACS numbers: 04.50.+h, 04.65.+e, 11.30.Qc

1. INTRODUCTION

A definitive answer to the question of how supersymmetry is broken in phenomenologically relevant theories of particle physics, be they supergravitational theories or superstrings, remains elusive. One candidate for a mechanism of supersymmetry breaking is the Polonyi model and its many generalizations [1-4]. This approach requires the explicit introduction of new chiral supermultiplets, in addition to those of the standard supersymmetric electroweak model and supergravity. Furthermore, the interactions of such fields with observable matter must be of a specific form; that is, they must form a "hidden" sector. Why such fields should exist, and how they come to be hidden from ordinary matter, is not well understood. A second candidate for a mechanism of supersymmetry breaking is non-perturbative, using gaugino condensation in a strongly interacting sector [5, 6]. Again, this mechanism requires extra superfields in the theory, in this case strongly interacting vector multiplets, and this sector must be sufficiently hidden from observable matter. Generically, it is not well understood why such a mechanism should occur, although in superstring theories this approach seems better motivated. In this paper, we will present what we believe to be a very different mechanism for supersymmetry breaking, which seems to overcome at least some of the shortcomings of the afore-mentioned approaches.

We begin by considering only the fields of the standard model and the graviton. In the usual way, we will assume that these fields are actually members of four- dimensional, $N = 1$ supermultiplets: chiral multiplets for quarks, leptons and Higgs fields, vector multiplets for

the strong and electroweak gauge fields, and the supergravitational multiplet for the graviton. These are the only superfields we will introduce, which distinguishes our approach from those discussed above. It is traditional, when generalizing the standard supersymmetric model to include supergravity, to admit higher-dimensional interaction terms composed of products of fields suppressed by the Planck mass, but to disallow higher-derivative interactions, be they purely gravitational or involving matter. Generically, there is no justification for this. Higher- derivative terms will appear at the same order when supermatter is coupled to supergravity. It is often said that such terms may indeed be present, but are irrelevant at low momentum since they will be suppressed by powers of momenta over the Planck mass. We wish to stress that this statement is naive. The reason is the following. It is well known that, in non-supersymmetric theories, higher- derivative terms in the equations of motion correspond to new degrees of freedom, in addition to the degrees of freedom of the original theory. This situation is amplified in supersymmetric theories where higher-derivative terms not only produce new bosonic and fermionic degrees of freedom, but also cause fields that were auxiliary, and, hence, not physical, in the original theory to become propagating [7]. We want to re-emphasize at this point that one is not introducing new superfields, rather the new degrees of freedom are arising out of derivatives of fields, or auxiliary fields, that have already been introduced. In two recent papers [8, 9], we have shown that these new degrees of freedom can have non-trivial potential energies possessing multiple vacua with vanishing cosmological constant. It is essential to realize that one must first identify the vacuum structure of these new degrees of freedom, and expand around one of the vacua, before taking the low- momentum limit. While it is true that the fluctuations around the vacuum will decouple at low momenta, the effect of the vacuum itself will not. It generically couples to the usual matter fields and survives to low momentum, much as the vacuum state of a grand-unified theory determines the gauge group structure of the low-energy effective theory. It is in this sense that the above statement is naive. This understood, it is important to ask what effect a non-trivial vacuum state can have on the low-energy effective matter Lagrangian. In a previous paper, we have shown that there exist vacua of the new degrees of freedom of pure higher-derivative supergravitation that spontaneously break supersymmetry. In this paper, we will couple such theories to the supersymmetric standard model and show that, at low momentum, such a non-trivial vacuum produces soft supersymmetry-breaking terms of fundamental phenomenological importance. Indeed, this mechanism seems sufficient to properly account for supersymmetry breaking in the standard supersymmetric model.

There is one further issue that we wish to touch upon before presenting our results. It is well known that higher-derivative theories tend to have at least some extra degrees of freedom whose propagation behavior is ghost-like. This seems to always be the case for matter fields [10] and is frequently the case in the gravitational sector [11-13]. A number of authors have pointed out that, at least in the context of gravity, this apparent ghost-like behavior may

be illusory [14-16]. This is a complicated issue that we would like to avoid in this paper. We have shown in a recent publication [9] that there is a consistent and interesting higher-derivative extension of supergravity that is, in fact, completely ghost-free. For simplicity, we will, in this paper, take this to be the higher-derivative extension of the standard model coupled to gravity. Hence, the theory discussed in this paper will be ghost-free. Why nature chooses this ghost-free extension of the standard model, or whether generic higher-derivative theories are actually consistent as well, will be left for future research.

2. Higher-Derivative Supergravity

In this section we will consider pure supergravitation only. Ordinary $N = 1$ Einstein supergravity is described by the Lagrangian

$$\mathcal{L} = -3M_P^2 \int d^4\theta E, \tag{2.1}$$

where E is the superdensity and M_P the Planck mass. It is well known that the minimal supergravity multiplet contains the graviton $e_m{}^a$, the gravitino $\psi_m{}^\alpha$, a complex scalar M and a real vector b_m as component fields. Expressed in terms of these component fields, Lagrangian (2.1) becomes [17]

$$e^{-1}\mathcal{L} = M_P^2 \left(-\tfrac{1}{2}\mathcal{R} - \tfrac{1}{3}MM^* + \tfrac{1}{3}b^m b_m\right) + \tfrac{1}{2}\epsilon^{klmn}\left(\bar{\psi}_k\bar{\sigma}_l\tilde{\mathcal{D}}_m\psi_n - \psi_k\sigma_l\tilde{\mathcal{D}}_m\bar{\psi}_n\right), \tag{2.2}$$

which describes bosonic Einstein gravitation coupled to a gravitino. The M and b_m fields are non-propagating auxiliary fields which can be eliminated from the Lagrangian using their equations of motion $M = 0$ and $b_m = 0$. It is important to note that it is not supersymmetry that is forcing M and b_m to be auxiliary, but rather it is the explicit choice of the Lagrangian (2.1). This Lagrangian was chosen because it describes the simplest supersymmetric extension of ordinary bosonic Einstein gravitation. However, were one to change it to include, for instance, higher-derivative gravitational terms, then it is no longer necessary that M and b_m be auxiliary fields. That is, they might begin to propagate and become physical. This is indeed the case, as was first noted at the linearized level in [7] and was established on the non-linear level in [18]. The same phenomenon in new minimal supergravity was discussed in [19, 20].

Recently [9], we constructed the most general class of ghost-free higher-derivative supergravity theories. The Lagrangian for such theories is given by

$$\mathcal{L} = -3M_P^2 \int d^4\theta E f\left(\frac{R}{m}, \frac{R^\dagger}{m}\right), \tag{2.3}$$

where f is an arbitrary dimensionless real function and we have chosen to normalize the scalar curvature superfield R by an arbitrary mass m. Expressed in terms of the component

fields, this Lagrangian becomes

$$e^{-1}\mathcal{L} = M_P^2 \Big\{ -\tfrac{1}{2}(f + Mf_M + M^*f_{M^*} - 4MM^*f_{MM^*} - 2b^m b_m f_{MM^*})\mathcal{R}$$

$$-\tfrac{3}{4}f_{MM^*}\mathcal{R}^2 + 3f_{MM^*}\partial^m M\partial_m M^* - 3f_{MM^*}(\nabla^m b_m)^2 + \cdots \Big\}, \quad (2.4)$$

where $f_{MM^*} = \partial^2 f/\partial M\partial M^*$ and we have dropped all fermionic terms and all bosonic terms inessential for this discussion. We see from this expression that Lagrangian (2.3) describes the supersymmetric extension of the higher-derivative bosonic $\mathcal{R} + \mathcal{R}^2$ theory. Furthermore, we see explicitly that both the complex M field and the longitudinal part of the b_m field are now propagating physical fields, as anticipated. It is well known that, in addition to the helicity-two graviton, $\mathcal{R} + \mathcal{R}^2$ bosonic gravitation contains a new propagating scalar field. Therefore, Lagrangian (2.3) contains four new bosonic scalar degrees of freedom in addition to the graviton. The fermionic superpartners of these new degrees of freedom can be identified in the higher-derivative fermionic terms that occur, but were not shown, in component Lagrangian (2.4). These new degrees of freedom must arrange themselves into two new chiral supermultiplets.

Written in this higher-derivative form, the physical content of this theory is obscure. However, it can be shown [18, 9] that by performing a super-Legendre transformation, Lagrangian (2.3) can be written in an equivalent form which consists of two chiral supermultiplets, denoted by Φ and Λ, coupled to ordinary Einstein supergravity with a particular Kähler potential and superpotential. Specifically, we showed in [9] that Lagrangian (2.3) is equivalent to the Lagrangian

$$\mathcal{L} = \int d^4\theta E \left\{ -3M_P^2 \exp\left(-\frac{1}{3}\frac{K}{M_P^2}\right) + \frac{W}{2R} + \frac{W^\dagger}{2R^\dagger} \right\}, \quad (2.5)$$

where

$$K = -3M_P^2 \ln\left\{ f\left(\frac{\Phi}{M_P}, \frac{\Phi^\dagger}{M_P}\right) + \frac{\Lambda + \Lambda^\dagger}{M_P} \right\},$$

$$W = 6m\Phi\Lambda. \quad (2.6)$$

Written in this form, it is relatively straightforward to determine the vacuum structure of the theory. Using the well-known formula for the potential energy [17], we find that

$$V = 12\left\{ f\left(\frac{A}{M_P}, \frac{A^*}{M_P}\right) + \frac{B + B^*}{M_P} \right\}^{-2} U(A, B), \quad (2.7)$$

where

$$U = M_P^2|A|^2\Big(f - 2f_A A - 2f_{A^*}A^* + 4f_{AA^*}|A|^2\Big) - f_{AA^*}^{-1}\Big|B - M_P f_A A + 2M_P f_{AA^*}|A|^2\Big|^2 \quad (2.8)$$

and A and B are the lowest components of chiral superfields Φ and Λ respectively. There is typically a local minimum of the potential at $\langle A \rangle = \langle B \rangle = 0$, with vanishing cosmological constant, where supersymmetry remains unbroken. We refer to this minimum as the trivial vacuum. Generically, however, we find that such potentials have other local minima. These non-trivial vacua can be arranged to have zero cosmological constant and, rather remarkably, are found to spontaneously break supersymmetry. It must be kept in mind that we are analyzing a theory of pure higher-derivative supergravitation. This theory involves the gravity supermultiplet only; no new fundamental supermultiplets have been introduced. What these results tell us is that the extra scalar degrees of freedom in such theories can have non-trivial vacua with vanishing cosmological constant which can spontaneously break supersymmetry. There is no need for a Polonyi-type field, gaugino condensates, or any other mechanism. This result seems sufficiently important that we will give a concrete example. Let us consider a function f of the form

$$f\left(\frac{R}{m}, \frac{R^\dagger}{m}\right) = 1 - 2\frac{RR^\dagger}{m^2} + \frac{1}{9}\frac{(RR^\dagger)^2}{m^4}. \tag{2.9}$$

We find that the associated potential energy has precisely two local minima, each with vanishing cosmological constant. The first is the trivial minimum at $\langle A \rangle = \langle B \rangle = 0$ which does not break supersymmetry. However, we find a second minimum, actually a ring of minima with zero cosmological constant, at

$$\begin{aligned}
\langle A \rangle &= M_P e^{i\theta}, \\
\langle B \rangle &= \tfrac{4}{3} M_P.
\end{aligned} \tag{2.10}$$

Computing the Kähler covariant derivatives at this vacuum, we find that

$$\begin{aligned}
\langle D_A W \rangle &= 32 m M_P, \\
\langle D_B W \rangle &= -\tfrac{15}{2} e^{i\theta} m M_P.
\end{aligned} \tag{2.11}$$

It follows that supersymmetry is spontaneously broken with strength $(mM_P)^{1/2}$. The associated gravitino mass is found to be

$$m_{3/2} = \tfrac{27}{8} m. \tag{2.12}$$

We now proceed to determine the effect of such a non-trivial vacuum on the low-energy effective Lagrangian of the standard model.

3. Higher-Derivative Supergravity Coupled to Matter

In this section, we will consider the matter superfields of the supersymmetric standard model and then couple these fields to higher-derivative supergravity. Since the couplings of the vector multiplets are essentially fixed by gauge invariance, we will here only consider the chiral supermultiplets. These we will generically label by Y_i, suppressing all gauge indices.

The supersymmetric standard model Lagrangian can be written in terms of a flat Kähler potential and a specific superpotential $g(Y_i)$ as

$$\mathcal{L}_{\text{Matter}} = \int d^4\theta \sum_i Y_i Y_i^\dagger + \left\{ \int d^2\theta g(Y_i) + \text{h.c.} \right\}. \tag{3.1}$$

The superpotential is trilinear in Y_i except for a possible Higgs bilinear term. In the following analysis, however, we do not need to use the exact form of g. It suffices to assume that it leads to a vacuum state where all $\langle Y_i \rangle$ either vanish, or are of the order of the electroweak scale. We will also take $\langle g \rangle = 0$ without loss of generality. Since supersymmetry is not observed at low-energy, the full model must also include explicit soft supersymmetry-breaking terms. The purpose of this section is to show that, when the standard model is coupled to higher-derivative supergravity, such soft terms are induced naturally when supersymmetry is spontaneously broken in the gravitational sector.

The supergravity extension of Lagrangian (3.1) is usually taken to be

$$\mathcal{L} = \int d^4\theta E \left\{ -3M_P^2 \exp\left(-\tfrac{1}{3} M_P^{-2} \sum_i Y_i Y_i^\dagger \right) + \frac{g}{2R} + \frac{g^\dagger}{2R^\dagger} \right\}. \tag{3.2}$$

In addition to Einstein supergravitation, Lagrangian (3.2) introduces a restricted set of higher-dimensional operators suppressed by powers of M_P. These operators involve products of the matter fields, but do not contain higher-derivative interactions. However, this is by no means the most general extension of Lagrangian (3.1) to the Planck scale. Generically, in addition to the Planck mass suppressed terms contained in (3.2), one can add higher-derivative terms involving both matter superfields and supergravity. There is a general class of such theories, where the higher-derivative terms are associated with the supergravity chiral superfield R, which are ghost-free. In this paper, we will restrict the discussion to theories of this type. A natural ghost-free higher-derivative extension of Lagrangian (3.1) is of the form

$$\mathcal{L} = \int d^4\theta E \left\{ -3M_P^2 F\left(\frac{R}{m}, \frac{R^\dagger}{m}, \frac{Y_i Y_i^\dagger}{M_P^2} \right) + \frac{g}{2R} + \frac{g^\dagger}{2R^\dagger} \right\}, \tag{3.3}$$

where F is a dimensionless real function and m is some mass-scale which does not have to be related to the Planck mass. We will assume, for the time being, that m is less than M_P by at least a few orders of magnitude. Further, we will assume that all the coefficients in the function F are of order unity, so that m and M_P are the only two mass scales in the theory.

Just as in the case of pure higher-derivative supergravity, we can make a super- Legendre transformation to put Lagrangian (3.3) into an equivalent second-order form. The transformed theory describes Einstein supergravity coupled to matter superfields Y_i plus two extra chiral supermultiplets, Φ and Λ. The superfields Φ and Λ represent the new degrees

of freedom arising from the higher-derivative supergravitational terms. Making the super-Legendre transformation gives the following Kähler potential and superpotential

$$K = -3M_P^2 \ln \left\{ F\left(\frac{\Phi}{M_P}, \frac{\Phi^\dagger}{M_P}, \frac{Y_i Y_i^\dagger}{M_P^2} \right) + \frac{\Lambda + \Lambda^\dagger}{M_P} \right\},$$

$$W = 6m\Phi\Lambda + g(Y_i).$$

(3.4)

The associated scalar potential energy is then found to be

$$V = 12 e^{2K/3M_P^2} \left(U_1 + U_2 + U_3 \right),$$

(3.5)

where

$$U_1 = m^2 AA^* \left(F - 2AF_A - 2A^* F_{A^*} + 4AA^* F_{AA^*} \right),$$

$$U_2 = -\frac{1}{2} \frac{m}{M_P} \left(A^* g + A g^* \right) - m^2 P_i (F_{ij^*})^{-1} P_j^*,$$

$$U_3 = -\frac{m^2}{M_P^2} \frac{\det F_{ij^*}}{\det X} \left| B - M_P (AF_A - 2AA^* F_{AA^*} - P_i (F_{ij^*})^{-1} F_{Aj^*}) \right|^2,$$

(3.6)

and we define

$$P_i = AF_i - 2AA^* F_{iA^*} - \frac{g_i}{6mM_P},$$

$$X = \begin{pmatrix} F_{AA^*} & F_{Aj^*} \\ F_{iA^*} & F_{ij^*} \end{pmatrix}.$$

(3.7)

As in the previous section, we write A and B for the lowest components for Φ and Λ respectively, and introduce y_i for the lowest component of Y_i, while $F_A = \partial F/\partial A$ and $F_i = \partial F/\partial y_i$. It is convenient to expand the function F as

$$F\left(\frac{\Phi}{M_P}, \frac{\Phi^\dagger}{M_P}, \frac{Y_i Y_i^\dagger}{M_P^2} \right) = f\left(\frac{\Phi}{M_P}, \frac{\Phi^\dagger}{M_P} \right) + \sum_i c_i \left(\frac{\Phi}{M_P}, \frac{\Phi^\dagger}{M_P} \right) \frac{Y_i Y_i^\dagger}{M_P^2} + \cdots$$

(3.8)

and insert this into potential energy (3.5). If we ignore, for the time being, all terms associated with matter superfields Y_i, then the potential energy reduces to the pure higher-derivative supergravity potential for the function f given in equation (2.7) of the previous section. There we demonstrated, for a wide class of functions f, that the potential energy has a stable, local minimum with non-vanishing vacuum expectation values $\langle A \rangle$ and $\langle B \rangle$ of the order of M_P. Furthermore, the cosmological constant can be chosen to be zero and, generically, supersymmetry is spontaneously broken with strength $(mM_P)^{1/2}$. We, henceforth, assume f is chosen to have these properties. Now let us restore the matter fields Y_i. We find that potential energy (3.5), although complicated, still possesses a vacuum with zero cosmological constant in which $\langle A \rangle$ and $\langle B \rangle$ are determined only by f and have the above properties. The y_i fields either have $\langle y_i \rangle = 0$, or non-zero values of the order of the electroweak scale, depending on the choice of the superpotential g.

We will now show that, at low-energy, the fluctuations around this vacuum reproduce the standard supersymmetric model described by Lagrangian (3.1) along with a specific set of soft supersymmetry-breaking terms. In this paper, we will consider the bosonic part of the Lagrangian only. We write $A = \langle A \rangle + a$ and $B = \langle B \rangle + b$ but, since $\langle y_i \rangle \ll M_P$, leave the y_i fields unexpanded. Inserting these expressions into the Lagrangian associated with (3.4) and dropping all terms suppressed by M_P^{-1} or more, we find that the low-energy theory separates into two pieces: a pure y_i part and a separate hidden sector depending on fields a and b, which have masses of order m, but do not interact with the y_i fields. We will, henceforth, ignore this hidden sector since, with the possible exception of cosmological consequences, it is physically irrelevant. The low-energy effective Lagrangian for the matter fields y_i is found to be

$$\mathcal{L}_{\text{eff}} = 3 e^{2\langle K \rangle / 3 M_P^2} \sum_i \langle c_i \rangle \, |\partial_m y_i|^2 - V_{\text{eff}}(y_i), \tag{3.9}$$

where

$$V_{\text{eff}}(y_i) = V_{\text{SUSY}}(y_i) + V_{\text{Soft}}(y_i). \tag{3.10}$$

The term V_{SUSY} is given by

$$V_{\text{SUSY}}(y_i) = -\tfrac{1}{3} e^{2\langle K \rangle / 3 M_P^2} \sum_i \frac{|g_i|^2}{\langle c_i \rangle}. \tag{3.11}$$

Provided that $\langle c_i \rangle < 0$, the kinetic energy term and V_{SUSY} can, after appropriate rescaling of the y_i fields, be written in the form of Lagrangian (3.1) with a modified superpotential g. Hence, these terms are supersymmetric. The remaining part, V_{Soft}, of the low- energy potential is found to be

$$V_{\text{Soft}}(y_i) = \sum_i m_{1,i}^2 |y_i|^2 + \left\{ m_2 g + \sum_i m_{3,i} y_i g_i + \text{c.c.} \right\}, \tag{3.12}$$

where

$$\begin{aligned}
m_{1,i}^2 &= 48 \frac{m^2}{M_P^2} \langle A A^* \rangle^2 \, e^{2\langle K \rangle / 3 M_P^2} \left\langle \frac{\partial^2 c_i}{\partial A \partial A^*} - c_i^{-1} \frac{\partial c_i}{\partial A} \frac{\partial c_i}{\partial A^*} \right\rangle, \\
m_2 &= -6 \frac{m}{M_P} \langle A^* \rangle \, e^{2\langle K \rangle / 3 M_P^2}, \\
m_{3,i} &= -2 \frac{m}{M_P} \langle A^* \rangle \, e^{2\langle K \rangle / 3 M_P^2} \left\{ 2 \left\langle A \frac{\partial \ln c_i}{\partial A} \right\rangle - 1 \right\}.
\end{aligned} \tag{3.13}$$

Unlike the previous case, V_{Soft} cannot be written in the form of Lagrangian (3.1) and, therefore, corresponds to genuine soft supersymmetry breaking. Expressions (3.12) and (3.13) are quite general, valid for a wide class of superpotentials g subject to the minor constraints

discussed previously. They simplify somewhat if we specialize to the case in which the matter superpotential is a homogeneous polynomial of degree n. We then have

$$\sum_i y_i g_i = ng. \tag{3.14}$$

This is the case in the R-parity invariant, supersymmetric standard model where the Higgs bilinear term is disallowed. Furthermore, since, in this model, the terms in g are trilinear, it follows that $n = 3$. Therefore, V_{Soft} in (3.12) simplifies to

$$V_{\text{Soft}} = \sum_i m_{1,i}^2 |y_i|^2 + \left\{ \sum_i m_{3,i}' y_i g_i + \text{c.c.} \right\}, \tag{3.15}$$

where

$$
\begin{aligned}
m_{1,i}^2 &= 48 \frac{m^2}{M_P^2} \langle AA^* \rangle^2 \, e^{2\langle K \rangle / 3M_P^2} \left\langle \frac{\partial^2 c_i}{\partial A \partial A^*} - c_i^{-1} \frac{\partial c_i}{\partial A} \frac{\partial c_i}{\partial A^*} \right\rangle, \\
m_{3,i}' &= -4 \frac{m}{M_P} e^{2\langle K \rangle / 3M_P} \left\langle AA^* \frac{\partial \ln c_i}{\partial A} \right\rangle.
\end{aligned}
\tag{3.16}
$$

We conclude that for a wide class of theories, including the supersymmetric standard model, spontaneous supersymmetry breaking in the higher-derivative supergravity sector is transmitted to the low-energy effective Lagrangian of observable matter as a specific set of soft, explicit supersymmetry-breaking operators. This result is reminiscent of the Polonyi mechanism with a general Kähler potential [2, 3], but, unlike that mechanism, it arises from pure supergravitation and does not require the introduction of a hidden supersymmetry-breaking sector.

The soft supersymmetry-breaking terms will lead, through radiative corrections, to the spontaneous breakdown of electroweak symmetry. If we assume, recalling $\langle A \rangle \simeq M_P$, that the $\langle c_i \rangle$-dependent coefficients of m or m^2 in (3.13) or (3.16) are of order unity then the scale of the electroweak symmetry breaking will be set by m. In this case, one must take $m \simeq 10^2 \text{GeV}$. The degree of fine-tuning required to do this is exactly the same as in the Polonyi models, and, hence, this is a completely viable approach to electroweak symmetry breaking. However, the higher-derivative supergravity theory described in this paper has a second, and quite novel, possibility for solving the hierarchy problem, in the context of the R-parity invariant, supersymmetric standard model. Note that the coefficients of the soft supersymmetry-breaking terms in (3.15) are explicit combinations of the vacuum expectation values of functions of c_i. Above, we assumed that these coefficients were of order unity. This then necessitated taking m to be of the order of the electroweak scale. Note, however, that if, for some reason, these combinations are small, of the order of 10^{-16} or so, then the parameter m can be chosen to be large. Thus, although the supersymmetry breaking in the matter sector is of the order of the electroweak scale, the gravitino mass and the masses of the hidden sector fields would be, generically, large. Consequently, in this case the so-called gravitino

and Polonyi cosmological problems [21] are automatically solved. Very small values for the $\langle c_i \rangle$-dependent coefficients are not as unexpected or unnatural as they may first appear to be. For example, suppose that

$$c_i \left(\frac{A}{M_P}, \frac{A^*}{M_P} \right) = \text{constant}, \tag{3.17}$$

for all values of i. It is immediately clear from (3.15) that V_{Soft} vanishes for any value of m. This is true despite the fact that supersymmetry is broken with strength $(mM_P)^{1/2}$ in the supergravity sector. This breaking is simply not transmitted to the low-energy observable theory. Condition (3.17) is very simple and it is intriguing to speculate whether it could arise naturally as a consequence of a symmetry of the higher-derivative theory. Now suppose the functions c_i are not strictly constant, but, instead, are slowly varying. It follows that V_{Soft} no longer vanishes. However, these terms can be very small, conceivably of electroweak strength, independent of the value of m. Such slowly varying functions could arise naturally, for example, from small, perhaps non- perturbative, breaking of the symmetry that enforced condition (3.17). This mechanism opens the possibility of theories involving the Planck mass only, with the gauge hierarchy arising from the structure of the functions c_i and not from the fine-tuning of the parameter m. The results of this section all involved approximations in which it was assumed that $m \ll M_P$. To explore this new mechanism, we will now consider the case where $m = M_P$.

4. COUPLING TO MATTER WITH $m = M_P$

We now repeat the analysis of the last section, but here we take $m = M_P$. It follows that we can no longer drop terms proportional to m/M_P. As a result, we must keep higher-order terms in the expansion of F. That is, instead of the expansion (3.8), we now consider

$$F \left(\frac{\Phi}{M_P}, \frac{\Phi^\dagger}{M_P}, \frac{Y_i Y_i^\dagger}{M_P^2} \right) = f \left(\frac{\Phi}{M_P}, \frac{\Phi^\dagger}{M_P} \right) + \sum_i c_i \left(\frac{\Phi}{M_P}, \frac{\Phi^\dagger}{M_P} \right) \frac{Y_i Y_i^\dagger}{M_P^2}$$
$$+ \sum_{ij} e_{ij} \left(\frac{\Phi}{M_P}, \frac{\Phi^\dagger}{M_P} \right) \frac{Y_i Y_i^\dagger}{M_P^2} \frac{Y_j Y_j^\dagger}{M_P^2} + \cdots . \tag{4.1}$$

For simplicity, we will consider only the R-parity invariant supersymmetric standard model with no Higgs bilinear term. Exactly as in the previous section, the theory has a vacuum state where $\langle A \rangle$ and $\langle B \rangle$ take the values set by the pure higher-derivative supergravity potential energy associated with the function f, while $\langle y_i \rangle$ are zero or of the order of the electroweak scale. Again, we assume that f is such that the vacuum spontaneously breaks supersymmetry with zero cosmological constant. However the scale of this breaking is now M_P. Furthermore, the masses of the fluctuations in A and B around the vacuum are now

generically also of the order of M_P and, hence, they completely decouple from the low-energy effective Lagrangian. We find that the low-energy Lagrangian is given by

$$\mathcal{L}_{\text{eff}} = 3e^{2\langle K \rangle/3M_P^2} \sum_i \langle c_i \rangle \, |\partial_m y_i|^2 - V_{\text{eff}}(y_i), \tag{4.2}$$

where

$$V_{\text{eff}} = V_{\text{SUSY}}(y_i) + V_{\text{Soft}}(y_i) + V_{\text{Quartic}}(y_i). \tag{4.3}$$

The terms V_{SUSY} and V_{Soft} have exactly the same form as in the previous section, but now with $m = M_P$. What is new here is that the potential energy also contains potentially hard supersymmetry-breaking operators of the form

$$V_{\text{Quartic}}(y_i) = \sum_{ij} \lambda_{ij} |y_i|^2 |y_j|^2, \tag{4.4}$$

where

$$\lambda_{ij} = \frac{12 \langle AA^* \rangle \, e^{2\langle K \rangle/3M_P^2}}{M_P^2} \left\langle \left| 1 - 2A \frac{\partial \ln c_i c_j}{\partial A} \right|^2 e_{ij} + 2 \left[1 - 2A^* \frac{\partial \ln c_i c_j}{\partial A^*} \right] A \frac{\partial e_{ij}}{\partial A} \right.$$
$$\left. + 2 \left[1 - 2A \frac{\partial \ln c_i c_j}{\partial A} \right] A^* \frac{\partial e_{ij}}{\partial A^*} + 4AA^* \frac{\partial^2 e_{ij}}{\partial A \partial A^*} \right\rangle$$
$$- \frac{e^{\langle K \rangle/3M_P^2}}{M_P^2} \left(\langle c_i \rangle \, m_{1,j}^2 + \langle c_j \rangle \, m_{1,i}^2 \right) - \frac{e^{-2\langle K \rangle/3M_P^2}}{48 M_P^2 \langle (AA^*)^2 f_{AA^*} \rangle} m_{1,i}^2 m_{1,j}^2. \tag{4.5}$$

In fact, a quartic term of this type, albeit proportional to m/M_P, was present in the previous section, as it is in the Polonyi models. Since in these theories $m \ll M_P$, such terms are automatically suppressed and, hence, are not discussed. In the present case, however, $m = M_P$ and these quartic terms are not automatically suppressed. In this paper, it will suffice to point out that the λ_{ij} coefficients can always be made to vanish or to be very small, for any choice of c_i, by appropriately adjusting the functions e_{ij} in equation (4.1). Once this is done, V_{Soft} can be made to be of the order of the electroweak scale by choosing the functions c_i to vary sufficiently slowly, exactly as was discussed in the previous section. To be more concrete, let us suppose that, for some reason,

$$c_i \left(\frac{A}{M_P}, \frac{A^*}{M_P} \right) = \text{constant}, \tag{4.6}$$

for all i and

$$e_{ij} \left(\frac{A}{M_P}, \frac{A^*}{M_P} \right) = \frac{h_{ij}(A/M_P) + (h_{ij}(A/M_P))^*}{|A/M_P|}, \tag{4.7}$$

for all i and j, where h_{ij} are arbitrary functions of A. Putting these expressions into (4.5) immediately yields

$$\lambda_{ij} = 0, \tag{4.8}$$

for all i and j. Therefore,the quartic terms exactly vanish. Furthermore, since $c_i = $ constant, we also have $V_{\text{Soft}} = 0$. In this case then, despite the fact that supersymmetry is spontaneously broken at order M_P in the supergravity sector, none of this breaking is transmitted to the low-energy effective Lagrangian. If the conditions (4.6) and (4.7) are now slightly altered, then the quartic terms will remain suppressed whereas an acceptable gauge hierarchy will be generated by the non-vanishing but small V_{Soft} terms. Conditions (4.6) and (4.7) are relatively simple and could arise naturally as the consequence of a symmetry of the higher-derivative theory. Small, perhaps non-perturbative, breaking of these symmetries would then give rise to the observed gauge hierarchy. This intriguing possibility will be discussed elsewhere [22]. We want to emphasize again that since, in these theories, both the gravitino and the supersymmetry-breaking scalar field masses are of order M_P, the gravitino and Polonyi cosmological problems [21] are solved. Furthermore, since the new degrees of freedom associated with higher-derivative supergravitation have masses of order M_P, this scenario could well arise within the context of superstrings [23].

Finally it is worth noting that the Lagrangian (3.3) is not the most general form of a ghost-free higher-derivative theory of supergravity coupled to matter. In general, one would not separate off the superpotential term $g/2R$ but, rather, include it in the function F, so that F becomes a function of Y_i and $Y_i^\dagger$ separately, giving

$$\mathcal{L} = -3M_P^2 \int d^4\theta E F(R, R^\dagger, Y_i, Y_i^\dagger). \tag{4.9}$$

As we will describe elsewhere [22], extending the analysis of this paper to the general case leads to essentially the same results, with only some minor modifications to the form of the low-energy effective theory.

Acknowledgments

References

[1] J. Poloyni, Budapest Preprint KFKI-93 (1977).

[2] S. K. Soni and H. A. Weldon, *Analysis of the supersymmetry breaking induced by $N = 1$ supergravity theories*, Phys. Lett. **126B** (1982), 215–219.

[3] L. Hall, J. Lykken, and S. Weinberg, *Supergravity as a messegener of supersymmetry breaking*, Phys. Rev. D **27** (1983), 2359–2378.

[4] H. P. Nilles, *Supersymmetry, supergravity and particle physics*, Phys. Rep **110** (1984), 1–162.

[5] S. Ferrara, L. Girardello, and H. P. Nilles, *Breakdown of local supersymmetry through gauge fermion condensates*, Phys. Lett. **125B** (1983), 457–460.

[6] D. Amati, K. Konishi, Y. Meurice, G. Rossi, and G. Veneziano, *Non-perturbative aspects in supersymmetric gauge theories*, Phys. Rep. **162** (1988), 169–248.

[7] S. Ferrara, M. T. Grisaru, and P. van Nieuwenhuizen, *Poincaré and conformal supergravity models with closed algebras*, Nucl. Phys. **B138** (1978), 430–444.

[8] A. Hindawi, B. A. Ovrut, and D. Waldram, *Two-dimensional higher-derivative supergravity and a new mechanism for supersymmetry breaking*, Nucl. Phys. **B471** (1996), 409–429, hep-th/9509174.

[9] A. Hindawi, B. A. Ovrut, and D. Waldram, *Four-dimensional higher-derivative supergravity and spontaneous supersymmetry breaking*, Nucl. Phys. **B476** (1996), 175–199, hep-th/9511223.

[10] A. Pais and G. E. Uhlenbeck, *On field theories with non-localized action*, Phys. Rev. **79** (1950), 145–165.

[11] K. Stelle, *Classical gravity with higher derivatives*, Gen. Rel. and Grav. **9** (1973), 353–371.

[12] A. Hindawi, B. A. Ovrut, and D. Waldram, *Consistent spin-two coupling and quadratic gravitation*, Phys. Rev. D **53** (1996), 5583–5596, hep-th/9509142.

[13] A. Hindawi, B. A. Ovrut, and D. Waldram, *Non-trivial vacua in higher-derivative gravitation*, Phys. Rev. D **53** (1996), 5597–5608, hep-th/9509147.

[14] E. T. Tomboulis, *Renormalizability and asymptotic freedom in quantum gravity*, Phys. Lett. **97B** (1980), 77–80.

[15] S. W. Hawking, *How's afraid of (higher-derivative) ghosts?*, Quantum Field Theory and Quantum Statistics (I. A. Batalin, C. S. Isham, and G. A. Vilkovisky, eds.), Adam Hilger, Bristol, 1987, pp. 129–139.

[16] E. T. Tomboulis, *Exact relation between Einstein and qudratic quantum gravity*, UCLA/95/TEP/39, hep-th/9601082.

[17] J. Wess and J. Bagger, *Supersymmetry and supergravity*, 2nd ed., Princeton University Press, Princeton, 1992.

[18] S. Cecotti, *Higher derivative supergravity is equivalent to standard supergravity coupled to matter*, Phys. Lett. **190B** (1987), 86–92.

[19] S. Cecotti, S. Ferrara, M. Porrati, and S. Sabharwal, *New minimal higher derivative supergravity coupled to matter*, Nucl. Phys. **B306** (1988), 160–180.

[20] B. A. Ovrut and S. Kalyana Rama, *Lorentz and $U(1)$ Chern-Simons terms in new minimal supergravity*, Phys. Lett. **254B** (1991), 132–138.

[21] G. D. Coughlan, N. Fischler, E. W. Kolb, S. Raby, and G. G. Ross, *Cosmological problems for the Polonyi potential*, Phys. Lett. **131B** (1983), 59–64.

[22] A. Hindawi, B. A. Ovrut, and D. Waldram, in preparation.

[23] K. Foerger, B. A. Ovrut, S. Theisen, and D. Waldram, *Higher derivative gravity in string theory*, Phys. Lett. **388B** (1996), 512–520, hep-th/9605145.